

# Novel duality in disorder driven local quantum criticality

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We find that competition between random Kondo and random magnetic correlations results in a quantum phase transition from a local Fermi liquid to a spin liquid. The local charge susceptibility turns out to have exactly the same critical exponent as the local spin susceptibility, suggesting novel duality between the Kondo singlet phase and the critical local moment state beyond the Landau-Ginzburg-Wilson symmetry breaking framework. This leads us to propose an enhanced symmetry at the local quantum critical point, described by an O(4) vector for spin and charge. The symmetry enhancement serves mechanism of electron fractionalization in critical impurity dynamics, where such fractionalized excitations are identified with topological excitations.

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To find mechanism of quantum number fractionalization has been an important direction of research in condensed matter physics [1]. An interesting proposal is spin fractionalization in the vicinity of a generic second order transition between symmetry unrelated orders, forbidden in the Landau-Ginzburg-Wilson (LGW) symmetry breaking framework [2]. An essential ingredient is that the original symmetry of a microscopic model is enhanced at the quantum critical point (QCP) [3], allowing the LGW forbidden duality between the symmetry unrelated orders [4]. The underlying mechanism is emergence of a topological term [2], originating from the enhanced symmetry at the QCP [3]. The topological phase assigns a non-trivial quantum number for the other phase to a topological soliton, the disorder parameter of one phase, and condensation of topological excitations results in the other phase, that is, emergence of duality. Such topological excitations are identified with fractionalized excitations at the deconfined QCP.

The generic second order transition between symmetry unrelated orders was also proposed in the context of heavy fermion quantum criticality, where a continuous transition seems to appear from an antiferromagnetic phase to a heavy fermion Fermi liquid [5]. Although there is no concrete construction for an effective field theory [6], an extended dynamical mean-field theory (DMFT) shows that the local spin susceptibility diverges at the same time as the antiferromagnetic spin susceptibility [7], implying an enhanced emergent symmetry because Kondo fluctuations are equally critical with antiferromagnetic ones.

In this letter we investigate the role of strong randomness in both Kondo and Ruderman-Kittel-Kasuya-Yosida (RKKY) correlations, where strong disorder leads us to the DMFT approximation naturally because disorder average covers spatial correlations [8]. We find that competition between random Kondo and random RKKY interactions results in a quantum phase transition from a local Fermi liquid to a spin liquid. The local charge suscepti-

bility turns out to have exactly the same critical exponent as the local spin susceptibility, suggesting novel duality between the Kondo singlet phase and the critical local moment state beyond the LGW framework. This leads us to propose an enhanced symmetry at the local QCP, described by an O(4) vector for spin and charge. We argue that the symmetry enhancement serves mechanism of impurity fractionalization at the local QCP, where spinons are identified with instantons in an O(4) nonlinear  $\sigma$  model type on a nontrivial manifold.

The previous works focused on the origin of non-Fermi liquid physics in the Kondo singlet phase without RKKY correlations [9]. The role of non-random RKKY interactions was examined in Ref. [10], where proximity of the Anderson localization for conduction electrons was proposed to be mechanism of the local Fermi liquid to spin liquid transition. We also point out the previous work on the disordered t-J model, where holes are doped into the spin liquid state of the disordered Heisenberg model, giving rise to marginal Fermi liquid phenomenology [11].

We start from an effective Anderson lattice model

$$H = \sum_{ij,\sigma} t_{ij} c_{i\sigma}^\dagger c_{j\sigma} + E_d \sum_{i\sigma} d_{i\sigma}^\dagger d_{i\sigma} + \sum_{ij} J_{ij} \mathbf{S}_i \cdot \mathbf{S}_j + \sum_{i\sigma} (V_i c_{i\sigma}^\dagger d_{i\sigma} + \text{H.c.}), \quad (1)$$

where  $t_{ij} = -\frac{t}{M\sqrt{z}}$  is a hopping integral for conduction electrons and

$$J_{ij} = \frac{J}{\sqrt{zM}} \varepsilon_i \varepsilon_j, \quad V_i = \frac{V}{\sqrt{M}} \varepsilon_i$$

are random RKKY and hybridization coupling constants, respectively. Here,  $M$  is the spin degeneracy. Randomness is given by

$$\overline{\varepsilon_i} = 0, \quad \overline{\varepsilon_i \varepsilon_j} = \delta_{ij}. \quad (2)$$

This model has two well known limits. In the  $V \rightarrow 0$  limit the random Heisenberg model results, where a spin

glass phase turns out to be unstable against a spin liquid state due to quantum fluctuations [12]. In the  $J \rightarrow 0$  limit the disordered Anderson lattice model was intensively investigated, as mentioned in the introduction, where the role of randomness in the energy level for localized electrons was revealed [9]. It is natural to expect a quantum phase transition from the Kondo singlet phase to the spin liquid state, increasing the ratio  $V/J$ .

The disorder average can be performed in the replica trick [8]. We observe that such an average neutralizes spatial correlations except the hopping term of conduction electrons. This leads us to the DMFT formulation [13]. Performing the DMFT approximation with the disorder average in the replica trick, we reach an effective local action for the strong random Anderson lattice model

$$\begin{aligned}
S_{dmft}^{replica} = & \int_0^\beta d\tau \left\{ \sum_{\sigma a} c_\sigma^{\dagger a}(\tau) (\partial_\tau - \mu) c_\sigma^a(\tau) \right. \\
& + \sum_{\sigma a} d_\sigma^{\dagger a}(\tau) (\partial_\tau + E_d) d_\sigma^a(\tau) \Big\} \\
& - \frac{V^2}{2M} \int_0^\beta d\tau \int_0^\beta d\tau' \sum_{\sigma\sigma'ab} [c_\sigma^{\dagger a}(\tau) d_\sigma^a(\tau) + d_\sigma^{\dagger a}(\tau) c_\sigma^a(\tau)] \\
& [c_{\sigma'}^{\dagger b}(\tau') d_{\sigma'}^b(\tau') + d_{\sigma'}^{\dagger b}(\tau') c_{\sigma'}^b(\tau')] \\
& - \frac{J^2}{2M} \int_0^\beta d\tau \int_0^\beta d\tau' \sum_{ab} \sum_{\alpha\beta\gamma\delta} S_{\alpha\beta}^a(\tau) R_{\beta\alpha\gamma\delta}^{ab}(\tau - \tau') S_{\delta\gamma}^b(\tau') \\
& + \frac{t^2}{M^2} \int_0^\beta d\tau \int_0^\beta d\tau' \sum_{ab\sigma} c_\sigma^{\dagger a}(\tau) G_{c\sigma\sigma}^{ab}(\tau - \tau') c_\sigma^b(\tau'), \quad (3)
\end{aligned}$$

where  $R_{\beta\alpha\gamma\delta}^{ab}(\tau - \tau')$  is the local spin-spin correlation function and  $G_{c\sigma\sigma}^{ab}(\tau - \tau')$  is the local electron propagator. Self-consistency is imposed in the Bethe lattice with an infinite number of lattice coordinations [14].

We solve the effective local action based on the U(1) slave-boson representation

$$d_\sigma^a = \hat{b}^{\dagger a} f_\sigma^a, \quad S_{\sigma\sigma'}^a = f_\sigma^{a\dagger} f_{\sigma'}^a - q_0^a \delta_{\sigma\sigma'} \quad (4)$$

with the single occupancy constraint  $b^{\dagger a}(\tau) b^a(\tau) + \sum_\sigma f_\sigma^{a\dagger}(\tau) f_\sigma^a(\tau) = 1$ , where  $q_0^a = \sum_\sigma f_\sigma^{a\dagger} f_\sigma^a / M$ . First, we show existence of a quantum phase transition based on the mean-field approximation, valid deep inside each stable phase. Second, we reveal the nature of the QCP beyond the mean-field approximation, where quantum corrections are introduced fully self-consistently, justified in the  $M \rightarrow \infty$  limit. In this study we are allowed to consider only paramagnetic and replica symmetric phases, protected due to strong quantum fluctuations of spin 1/2, consistent with the previous studies [12].

In the slave-boson mean-field approximation we replace the holon operator with its expectation value  $\langle \hat{b}^a \rangle \equiv b^a$ . Then, we reach self-consistent equations for self-energy corrections in the replica symmetric phase

$$\Sigma_c(i\omega_l) = \frac{V^2}{M} G_f(i\omega_l) |b|^2 + \frac{t^2}{M^2} G_c(i\omega_l),$$

$$\begin{aligned}
\Sigma_f(i\omega_l) &= \frac{V^2}{M} G_c(i\omega_l) |b|^2 + \frac{J^2}{2M} T \sum_s \sum_{\nu_m} G_f(i\omega_l - i\nu_m) \\
& [R_{ss\sigma\sigma}(i\nu_m) + R_{\sigma ss\sigma}(-i\nu_m)], \\
\Sigma_{cf}(i\omega_l) &= \frac{V^2}{M} G_{fc}(i\omega_l) (b^2)^* - n \frac{V^2}{M} (b^2)^* \sum_s \langle f_s^\dagger c_s + c_s^\dagger f_s \rangle, \\
\Sigma_{fc}(i\omega_l) &= \frac{V^2}{M} G_{cf}(i\omega_l) b^2 - n \frac{V^2}{M} b^2 \sum_s \langle f_s^\dagger c_s + c_s^\dagger f_s \rangle, \\
R_{\sigma ss\sigma}(i\nu_m) &= -\frac{1}{\beta} \sum_{\omega_l} G_{f\sigma}(i\nu_m + i\omega_l) G_{fs}(i\omega_l), \quad (5)
\end{aligned}$$

where the Green's functions are given by

$$\begin{aligned}
& \begin{pmatrix} G_c(i\omega_l) & G_{fc}(i\omega_l) \\ G_{cf}(i\omega_l) & G_f(i\omega_l) \end{pmatrix} \\
& = \begin{pmatrix} i\omega_l + \mu - \Sigma_c(i\omega_l) & -\Sigma_{cf}(i\omega_l) \\ -\Sigma_{fc}(i\omega_l) & i\omega_l - E_d - \lambda - \Sigma_f(i\omega_l) \end{pmatrix}^{-1}.
\end{aligned}$$

$n$  is the replica index, set to be zero. Self-consistent equations for holon condensation and an effective chemical potential  $\lambda$  are

$$\begin{aligned}
& b \left[ \lambda + 2V^2 T \sum_{\omega_l} G_c(i\omega_l) G_f(i\omega_l) \right. \\
& + V^2 T \sum_{\omega_l} \left\{ G_{fc}(i\omega_l) G_{fc}(i\omega_l) + G_{cf}(i\omega_l) G_{cf}(i\omega_l) \right\} \Big] = 0, \\
& |b|^2 + \sum_\sigma \langle f_\sigma^\dagger f_\sigma \rangle = 1. \quad (6)
\end{aligned}$$

The main difference between the clean and disordered cases is that the off diagonal Green's function  $G_{fc}(i\omega_l)$  should vanish in the presence of randomness in  $V$  with its zero mean value while it is proportional to the condensation  $b$  when the average value of  $V$  is finite [15]. In the present situation we find  $b^a = \langle f_\sigma^{a\dagger} c_\sigma^a \rangle = 0$  while  $(b^a)^* b^b = \langle f_\sigma^{a\dagger} c_\sigma^a c_\sigma^{b\dagger} f_\sigma^b \rangle \equiv |b|^2 \delta_{ab} \neq 0$ . This implies that the Kondo singlet phase is not characterized by the holon condensation but described by finite density of holons. It is important to notice that this gauge invariant order parameter does not cause any kinds of symmetry breaking for the Kondo effect. This cures the artificial finite temperature transition in the slave-boson mean-field theory of the Anderson lattice model without randomness [15].

Figure 1 shows the phase diagram in the plane of  $(V, J)$ , where  $V$  and  $J$  are variances for the Kondo and RKKY interactions, respectively. The phase boundary is characterized by  $|b|^2 = 0$ , below which  $|b|^2 \neq 0$  appears to cause effective hybridization between conduction electrons and localized fermions. In the left panel of Fig. 2 one finds that the effective hybridization enhances the scattering rate of conduction electrons dramatically around the Fermi energy while the scattering rate for localized electrons becomes reduced at the resonance energy. This self-energy effect reflects the spectral function in the right panel of Fig. 2, where the pseudogap feature arises in conduction electrons while the

sharply defined peak appears in localized electrons, identified with the Kondo resonance although the description of the Kondo effect differs from the clean case. Increasing the RKKY coupling, the Kondo effect is suppressed as expected. In the Kondo singlet phase the local spin susceptibility shows the typical  $\omega$ -linear behavior in the low frequency limit, nothing but the Fermi liquid physics for spin correlations. Increasing  $J$ , incoherent spin correlations are enhanced, consistent with spin liquid physics.

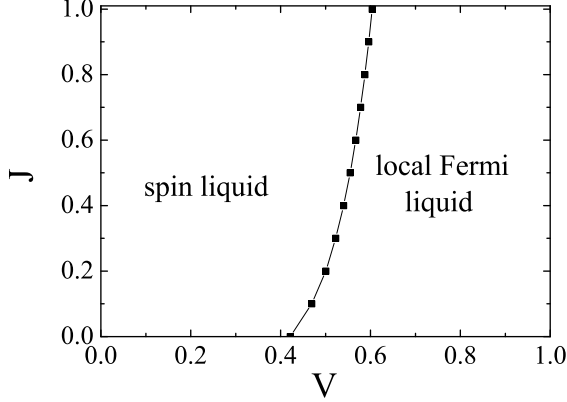


FIG. 1: The phase diagram of the strongly disordered Anderson lattice model in the DMFT approximation ( $E_d = -1$ ,  $\mu = 0$ ,  $T = 0.01$ ,  $t = 1$ ,  $M = 2$ ).

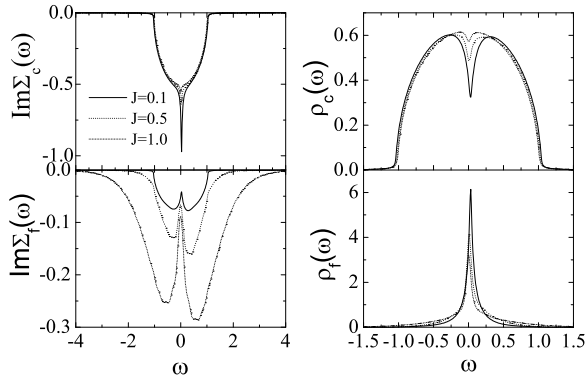


FIG. 2: Left: The imaginary part of the self-energy of conduction electrons and that of localized electrons for various values of  $J$ . Right: Density of states of conduction ( $\rho_c(\omega)$ ) and localized ( $\rho_f(\omega)$ ) electrons for various values of  $J$ . We used  $V = 0.5$ ,  $E_d = -0.7$ ,  $\mu = 0$ ,  $T = 0.01$ ,  $t = 1$ , and  $M = 2$ .

The nature of the local QCP is uncovered in the non-crossing approximation, exact in the  $M \rightarrow \infty$  limit [16].

Self-consistent equations for self-energy corrections are

$$\begin{aligned}\Sigma_c(\tau) &= \frac{V^2}{M} G_f(\tau) G_b(-\tau) + \frac{t^2}{M^2} G_c(\tau), \\ \Sigma_f(\tau) &= \frac{V^2}{M} G_c(\tau) G_b(\tau) - J^2 [G_f(\tau)]^2 G_f(-\tau), \\ \Sigma_b(\tau) &= V^2 G_c(-\tau) G_f(\tau),\end{aligned}\quad (7)$$

where the holon propagator is  $G_b(i\nu_l) = \left( i\nu_l - \lambda - \Sigma_b(i\nu_l) \right)^{-1}$ . When the second terms are neglected in the first and second equations, these are reduced to those of the multi-channel Kondo effect [16]. Power-law solutions are well known in the regime of  $1/T_K \ll \tau \ll \beta = 1/T$ , where  $T_K = D[\Gamma_c/\pi D]^{1/M} \exp[\pi E_d/M\Gamma_c]$  is an effective Kondo temperature [17] with the conduction bandwidth  $D$  and effective hybridization  $\Gamma_c = \pi \rho_c \frac{V^2}{M}$ . The RKKY interaction will reduce the effective hybridization, where  $\Gamma_c$  is replaced with  $\Gamma_c^J \approx \pi \rho_c (\frac{V^2}{M} - J^2)$ .

Our power-law ansatz is as follows

$$\begin{aligned}G_c(\tau) &= A_c \beta^{-\Delta_c} g_c\left(\frac{\tau}{\beta}\right), \quad G_f(\tau) = A_f \beta^{-\Delta_f} g_f\left(\frac{\tau}{\beta}\right), \\ G_b(\tau) &= A_b \beta^{-\Delta_b} g_b\left(\frac{\tau}{\beta}\right),\end{aligned}\quad (8)$$

where  $g_\alpha(x) = \left( \frac{\pi}{\sin(\pi x)} \right)^{\Delta_\alpha}$  is the scaling function [17] at finite temperatures with  $\alpha = c, f, b$ .  $A_c$ ,  $A_f$ , and  $A_b$  are numerical constants. Inserting these expressions into Eq. (7), we find two fixed point solutions. One coincides with the multi-channel Kondo effect, given by  $\Delta_c = 1$ , and  $\Delta_f = \frac{M}{M+1}$ ,  $\Delta_b = \frac{1}{M+1}$  with  $M = 2$ , where contributions from spin fluctuations to self-energy corrections are irrelevant, compared with holon fluctuations. The other is  $\Delta_c = 1$  and  $\Delta_f = \Delta_b = \frac{1}{2}$ , where spin correlations are critical as much as holon fluctuations. One can understand the critical exponent  $\Delta_f = 1/2$  as the proximity of the spin liquid physics [12]. Actually, we find the local spin susceptibility  $\Im\chi(\omega) = A_f^2 \left( \frac{1}{T} \right)^{1-2\Delta_f} \Phi\left(\frac{\omega}{T}\right)$ , where the scaling function is

$$\Phi(x) = 2(2\pi)^{2\Delta_f-1} \sinh\left(\frac{x}{2}\right) \frac{\Gamma\left(\Delta_f + i\frac{x}{2\pi}\right) \Gamma\left(\Delta_f - i\frac{x}{2\pi}\right)}{\Gamma(2\Delta_f)},$$

which coincides with the spin spectrum of the spin liquid state when  $V = 0$ . In this respect the second fixed point is the genuine critical solution for the local Fermi liquid to spin liquid transition.

It is straightforward to see that the critical exponent of the local spin susceptibility is exactly the same as that of the local charge susceptibility ( $2\Delta_f = 2\Delta_b = 1$ ), proportional to  $1/\tau$ . Such an unexpected scaling behavior proposes an enhanced symmetry, allowing us to construct an effective local field theory in terms of an  $O(4)$  vector

$$\Psi^a(\tau) = \begin{pmatrix} S^a(\tau) \\ \rho^a(\tau) \end{pmatrix},$$

$$Z_{eff} = \int D\Psi^a(\tau) \delta(|\Psi^a(\tau)|^2 - 1) e^{-S_{eff}},$$

$$S_{eff} = -\frac{g^2}{2M} \int_0^\beta d\tau \int_0^\beta d\tau' \Psi^{aT}(\tau) \Upsilon^{ab}(\tau - \tau') \Psi^b(\tau') + S_{top}, \quad (9)$$

where  $\Upsilon^{ab}(\tau - \tau')$  determines dynamics of the O(4) vector, resulting from spinon and holon dynamics in principle.  $g \propto V/J$  is an effective coupling constant, and  $S_{top}$  is a possible topological term.

One can represent the O(4) vector generally as follows

$$\Psi^a : \tau \longrightarrow \begin{pmatrix} \sin \theta^a(\tau) \sin \phi^a(\tau) \cos \varphi^a(\tau), \\ \sin \theta^a(\tau) \sin \phi^a(\tau) \sin \varphi^a(\tau), \\ \sin \theta^a(\tau) \cos \phi^a(\tau), \cos \theta^a(\tau) \end{pmatrix},$$

where  $\theta^a(\tau), \phi^a(\tau), \varphi^a(\tau)$  are three angle coordinates for the O(4) vector. It is essential to observe that the target manifold for the O(4) vector is not a simple sphere type, but more complicated because the last component of the O(4) vector is the charge density field. Its positiveness results in a periodicity, given by  $\Psi^a(\theta^a, \phi^a, \varphi^a) = \Psi^a(\pi - \theta^a, \phi^a, \varphi^a)$ . This folded space structure allows a nontrivial topological excitation. Suppose the boundary configuration of  $\Psi^a(0, \phi^a, \varphi^a; \tau = 0)$  and  $\Psi^a(\pi, \phi^a, \varphi^a; \tau = \beta)$ , connected by  $\Psi^a(\pi/2, \phi^a, \varphi^a; 0 < \tau < \beta)$ . Interestingly, this configuration is *topologically* distinguishable from the configuration of  $\Psi^a(0, \phi^a, \varphi^a; \tau = 0)$  and  $\Psi^a(0, \phi^a, \varphi^a; \tau = \beta)$  with  $\Psi^a(\pi/2, \phi^a, \varphi^a; 0 < \tau < \beta)$  because of the folded structure. This topological excitation carries a spin quantum number 1/2 in its core, given by  $\Psi^a(\pi/2, \phi^a, \varphi^a; 0 < \tau < \beta) = (\sin \phi^a(\tau) \cos \varphi^a(\tau), \sin \phi^a(\tau) \sin \varphi^a(\tau), \cos \phi^a(\tau), 0)$ .

This is the spinon excitation, described by an O(3) nonlinear  $\sigma$  model with the nontrivial spin correlation function  $\Upsilon^{ab}(\tau - \tau')$ , where the topological term is reduced to the single spin Berry phase term in the instanton core [3].

In this local impurity picture the local Fermi liquid phase is described by gapping of instantons while the spin liquid state is characterized by condensation of instantons. Of course, the low dimensionality does not allow condensation, resulting in critical dynamics for spinons. This scenario clarifies the LGW forbidden duality between the Kondo singlet and the critical local moment for the impurity state, allowed by the presence of the topological term.

We explicitly checked that the similar result can be found in the extended DMFT for the clean Kondo lattice model, where two fixed point solutions are allowed [18].

One is the same as the multi-channel Kondo effect and the other is essentially the same as the second solution in this paper. In this respect we believe that the present scenario works in the extended DMFT framework although applicable to only two spatial dimensions [7].

One may suspect the applicability of the DMFT framework for this disorder problem. However, the hybridization term turns out to be exactly local in the case of strong randomness while the RKKY term is safely approximated to be local for the spin liquid state, expected to be stable against the spin glass phase in the quantum spin case [13]. This situation should be distinguished from the clean case, where the DMFT approximation causes several problems such as the stability of the spin liquid state and strong dependence of the dimension of spin dynamics [7, 11, 18].

In conclusion, we proposed novel duality between the Kondo and critical local moment phases in the strongly disordered Anderson lattice model. This duality serves mechanism of impurity fractionalization at the local QCP, where spinons are identified with instantons in an O(4) nonlinear  $\sigma$  model on a nontrivial manifold.

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